Properties of collective flow and pion production in intermediate-energy heavy-ion collisions with a relativistic quantum molecular dynamics model*

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The relativistic mean-field approach was implemented in the Lanzhou quantum molecular dynamics transport model (LQMD.RMF). Using the LQMD.RMF, the properties of collective flow and pion production were investigated systematically for nuclear reactions with various isospin asymmetries. The directed and elliptic flows of the LQMD.RMF are able to describe the experimental data of STAR Collaboration. The directed flow difference between free neutrons and protons was associated with the stiffness of the symmetry energy, that is, a softer symmetry energy led to a larger flow difference. For various collision energies, the ratio between the π^- and π^+ yields increased with a decrease in the slope parameter of the symmetry energy. When the collision energy was 270 MeV/nucleon, the single ratio of the pion transverse momentum spectra also increased with decreasing slope parameter of the symmetry energy in both nearly symmetric and neutron-rich systems. However, it remained difficult to determine the dependence of the double ratio on the stiffness of the symmetry energy.

Keywords: Heavy-ion collision, Collective flow, Pion production, Symmetry energy, Relativistic mean field

I. INTRODUCTION

The equation of state (EOS) of nuclear matter, which originates from the interactions between nuclear matter, plays an important role in nuclear physics and astrophysics. Heavy-ion collisions, the properties of nuclei, and neutron stars (NSs) have been widely studied to extract the nuclear EOS. Because nuclear many-body problems are highly nonlinear and the EOS is not a directly observable quantity in experiments, there are still some uncertainties in the EOS despite great efforts [1–6]. For instance, the EOS extracted from the GW170817 event has uncertainties at high nuclear densities [2]. Although the EOS can be extracted from the properties of NSs, the internal composition of NSs is still poorly understood. The core of an NS may contain exotic materials such as hyperons, kaons, pions, and deconfined quark matter [7–12]. Heavy-ion collisions in terrestrial laboratories provide a unique opportunity to study both the EOS and exotic materials.

Collective flows of heavy-ion collisions were proposed in the 1970s and first detected in experiments at Bevalac [13–16]. Because collective flows are associated with nucleon-nucleon interactions, nucleon-nucleon scattering, etc., collective flows have been used to extract the nuclear EOS [1]. Collective flows are also helpful for understanding the phase transition between hadronic and quark matter. Generally, when a phase transition between hadronic and quark matter occurs, the collective flows of heavy-ion collisions indicate a soft EOS [17–21]. In addition, the ratios of the isospin particles in heavy-ion collisions, such as π^-/π^+, K^0/K^+, and Σ^-/Σ^+, are thought to be sensitive to the stiffness of the EOS [22–28]. The production of pions and kaons has been experimentally measured in ^{197}Au+^{197}Au collisions. The K^+ production predicted by various transport models favors a soft EOS of isospin-symmetric nuclear matter at high baryon densities [29–33]. The π^-/π^+ ratio predicted by various transport models is still model-dependent [34–37]. Based on the FOPI data for the π^-/π^+ ratio [38], some results favor a stiff symmetry energy (isospin asymmetric part of the EOS) [34, 35], whereas others imply a soft symmetry energy [36, 37]. Recently, by analyzing the ratios of charged pions in ^{132}Sn $+^{124}$ Sn, ^{112}Sn $+^{124}$ Sn, and ^{108}Sn $+^{112}$ Sn collisions [39], a slope of the symmetry energy ranging from 42 to 117 MeV was suggested [40, 41]. The collective flows and ratios of charged pions are still worth studying to find the sources of the difference in various transport models and to extract information about the EOS from heavy-ion collisions.

Quantum molecular dynamics (QMD) is a popular transport model that has been developed into many versions and used to successfully describe the properties of nuclear matter, nuclei, mesons, and hyperons [33, 34, 42–55]. In high-energy heavy-ion collisions, the relativistic effects should be considered in QMD because they become significant. The RQMD approach has been proposed for this purpose [42, 43]. Recently, relativistic mean field theory (RMF) was implemented in RQMD (RQMD.RMF) [44–46]. The RQMD.RMF has been successfully applied to investigate the collective flows of hadrons [44–46]. In this study, we implemented RMF theory with isovector-vector and isovector-scalar fields in the Lanzhou quantum molecular dynamics model (LQMD.RMF). The channels for the generation and decay of resonances (Δ(1232), N*(1440), N*(1535), etc.), hyperons, and mesons were included in the previous LQMD model [33, 34, 47–50]. Using the LQMD.RMF, we explored the relationship between the symmetry energy and the properties of the collective flow and pion production.

The remainder of this paper is organized as follows. In Sec. II, we briefly introduce the formulas and approaches used in this study. The formulas include RMF theory, the dispersion relation, and the production of pions. The results and discussion are presented in Sec. III. Finally, a summary is presented in Sec. IV.

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* This study was supported by the National Natural Science Foundation of China (Nos. 12147106, 12175072, and 11722546) and the Talent Program of South China University of Technology (Project No. 20210115).

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II. FORMALISM

A. Relativistic mean field approach

The RMF interaction is achieved by exchanging mesons. Scalar and vector mesons provide medium-range attraction and short-range repulsion between nucleons, respectively [56]. The nonlinear self-interaction of the σ meson is introduced to reduce the incompressibility of a reasonable domain [57]. To investigate the properties of symmetry energy, we also consider the isovector-vector ρ [58] and isovector-scalar δ mesons [59]. The Lagrangian density is expressed as [59, 60]

$$\begin{aligned}\mathcal{L} = & \bar{\psi}[\gamma_\mu(i\partial^\mu - g_\omega\omega^\mu - g_\rho\vec{\tau}\cdot\vec{b}^\mu) - (M_N - g_\sigma\sigma \\ & -g_\delta\vec{\tau}\cdot\vec{\delta})]\psi + \frac{1}{2}(\partial_\mu\sigma\partial^\mu\sigma - m_\sigma^2\sigma^2) - \frac{1}{3}g_2\sigma^3 \\ & -\frac{1}{4}g_3\sigma^4 + \frac{1}{2}m_\omega^2\omega_\mu\omega^\mu - \frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \frac{1}{2}m_\rho^2\vec{b}_\mu\vec{b}^\mu \\ & -\frac{1}{4}\vec{B}_{\mu\nu}\vec{B}^{\mu\nu} + \frac{1}{2}(\partial_\mu\vec{\delta}\cdot\partial^\mu\vec{\delta} - m_\delta^2\vec{\delta}^2), \end{aligned} \tag{1}$$

where $M_N = 938$ MeV is the nucleon mass in free space, g_i with $i = \sigma, \omega, \rho, \delta$ is the coupling constant between nucleons, m_i with $i = \sigma, \omega, \rho, \delta$ denotes the meson mass, g_2 and g_3 are the coupling constants for the nonlinear self-interaction of the σ meson, and $F_{\mu\nu} = \partial_\mu\omega_\nu - \partial_\nu\omega_\mu$ and $\vec{B}_{\mu\nu} = \partial_\mu\vec{b}_\nu - \partial_\nu\vec{b}_\mu$ are the strength tensors of the ω and ρ mesons, respectively. The equations of motion for the nucleon and meson are obtained from the Euler–Lagrange equations and are written as:

$$[i\gamma^\mu\partial_\mu - g_\omega\gamma^0\omega_0 - g_\rho\gamma^0 b_0\tau_3 - (M_N - g_\sigma\sigma - g_\delta\tau_3\delta_3)]\psi = 0 \tag{2}$$

$$m_\sigma^2\sigma + g_2\sigma^2 + g_3\sigma^3 = g_\sigma\bar{\psi}\psi = g_\sigma\rho_S \tag{3}$$

$$m_\omega^2\omega_0 = g_\omega\bar{\psi}\gamma^0\psi = g_\omega\rho \tag{4}$$

$$m_\rho^2 b_0 = g_\rho\bar{\psi}\gamma^0\tau_3\psi = g_\rho\rho_3, \tag{5}$$

$$m_\delta^2\delta_3 = g_\delta\bar{\psi}\tau_3\psi = g_\delta\rho_{S3}, \tag{6}$$

where ρ and ρ_S are the baryon and scalar densities, respectively, $\rho_3 = \rho_p - \rho_n$ is the difference between the proton and neutron densities, and $\rho_{S3} = \rho_{Sp} - \rho_{Sn}$ is the difference between the proton and neutron scalar densities.

In the RMF approximation, the energy density is given by

$$\begin{aligned}\epsilon = & \sum_{i=n,p} 2\int_0^{p_F}\frac{d^3p}{(2\pi)^3}\sqrt{p^2 + M_i^{*2}} + \frac{1}{2}m_\sigma^2\sigma^2 + \frac{1}{3}g_2\sigma^3 \\ & + \frac{1}{4}g_3\sigma^4 + \frac{1}{2}m_\omega^2\omega_0^2 + \frac{1}{2}m_\rho^2 b_0^2 + \frac{1}{2}m_\delta^2\delta_3^2, \end{aligned} \tag{7}$$

where p_F denotes the nucleon Fermi momentum, and $M_i^* = M_N - g_\sigma\sigma \mp g_\delta\delta_3$ ($-$ proton, $+$ neutron) denotes the effective nucleon mass. The isospin splitting of the effective nucleon mass $M_n^* - M_p^*$ still has a large uncertainty at this point. Analyses of nucleon-nucleus scattering data based on the optical model favor $M_n^* - M_p^* > 0$ [61, 62]. Calculations based on Brueckner theory [63–65] and density-dependent relativistic Hartree-Fock [66] also indicate $M_n^* - M_p^* > 0$. However, $M_n^* - M_p^* < 0$ is predicted by the transport model for heavy-ion collisions [67, 68] and nonlinear RMF models [59, 60, 69]. In addition, both $M_n^* - M_p^* < 0$ and $M_n^* - M_p^* > 0$ can be predicted by the point-coupling RMF [69] and Skyrme and Gogny forces [70–76]. Because the Lagrangian density in this study is the same as that in Ref. [59] and [60] except for the parameter settings, as shown in Fig. 1, the negative isospin splitting of the effective nucleon mass $M_n^* - M_p^* < 0$ is consistent with that in Ref. [59] and [60]. In the nonlinear RMF model, the isospin splitting of the effective nucleon mass is primarily determined by the coupling strength of the δ meson. The large coupling strength of the δ meson results in large isospin splitting of the effective nucleon mass. When the coupling strength of the δ meson is zero (set1), there is no isospin splitting of the effective nucleon mass.

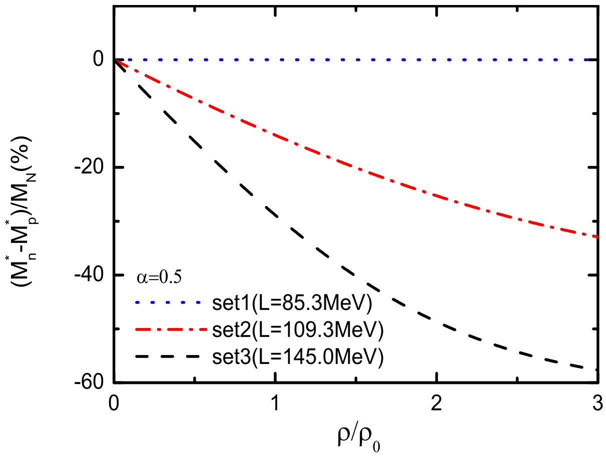

Fig. 1. (Color online) Difference between neutron and proton effective masses as a function of the baryon density.

Using the isospin asymmetry parameter $\alpha = (\rho_n - \rho_p)/(\rho_n + \rho_p)$, the symmetry energy can be written as [59]

$$\begin{aligned}E_{sym} = & \frac{1}{2}\frac{\partial^2 E(\rho,\alpha)}{\partial\alpha^2}|_{\alpha=0} = \frac{1}{2}\frac{\partial^2[\epsilon(\rho,\alpha)/\rho]}{\partial\alpha^2}|_{\alpha=0} \\ = & \frac{1}{6}\frac{p_F^2}{E_F^*} + \frac{1}{2}f_\rho\rho - \frac{f_\delta}{2}\frac{M^{*2}\rho}{E_F^{*2}[1 + f_\delta A(p_F, M^*)]}, \end{aligned} \tag{8}$$

where $f_i \equiv \frac{g_i^2}{m_i^2}$, $i = \rho, \delta$, and $E_F^* = \sqrt{p_F^2 + M^{*2}}$ and $M^* = M_N - g_\sigma\sigma$ are the effective nucleon masses of the symmetric

nuclear matter. The integral $A(p_F, M^*)$ is defined as

$$A(p_F, M^*) = \frac{4}{(2\pi)^3}\int d^3p \frac{p^2}{(p^2+M^{*2})^{3/2}}$$
$$= 3(\frac{\rho_S}{M^*} - \frac{\rho}{E_F^*}). \quad (9)$$

Based on the symmetry energy E_{sym}, the slope L and curvature K_{sym} of the symmetry energy at the saturation density ρ_0 can be obtained as

$$L = 3\rho_0(\frac{\partial E_{sym}}{\partial \rho})|_{\rho=\rho 0}$$
$$K_{sym} = 9\rho_0^2(\frac{\partial^2 E_{sym}}{\partial \rho^2})|_{\rho=\rho 0}. \quad (10)$$

In this study, we set the saturation density to $\rho_0 = 0.16 fm^{-3}$, and the binding energy per particle of the symmetric nuclear matter was set to $E/A - M_N = -16$ MeV. For symmetric nuclear matter, we set set1, set2, and set3 models to be the same as the result of vanishing isospin asymmetry. As shown in Table 1 and Fig. 2, the symmetry energy of set1, set2, and set3 was set to be 31.6 MeV at the saturation density. Set1 contained only ρ mesons; however, set2 and set3 contained both the ρ and δ mesons. For set1, when the symmetry energy was set to be 31.6 MeV at the saturation density, the coupling parameter g_ρ was fixed and the slope of symmetry was fixed at $L = 85.3$ MeV. For set2 and set3, when the symmetry energy was fixed at 31.6 MeV, the slope of the symmetry energy was obtained by varying the coupling parameters g_ρ and g_δ. Because the effective mass M^* of the above models for symmetric nuclear matter was the same, the symmetry energy with both the ρ and δ mesons could not be softer than that of set1 containing only ρ mesons. To broaden the range of the slope parameters, we set the slope parameter of set2 and set3 to be 109.3 and 145.0 MeV by varying the coupling parameters g_ρ and g_δ, respectively. A broader range of slope parameters would be helpful for understanding the relationship between the properties of symmetry energy and the observables of heavy-ion collisions. The curvature of the symmetry energy $K_{\rm sym}$, which is a higher-order expansion coefficient of the symmetry energy compared to the slope parameter L, may also affect the properties of the symmetry energy and the observables of heavy-ion collisions at densities far beyond the saturation density. K_{sym} of set1, set2, and set3 is obtained as -15, 141, and 391 MeV, respectively.

B. Relativistic quantum molecular dynamics approach

To investigate high-energy heavy-ion collisions, RQMD was proposed [42, 43]. Recently, RMF was implemented in RQMD [44–46]. In RQMD, for an N-body system, there are $4N$ position coordinates q_i^μ and $4N$ momentum coordinates p_i^μ $(i = 1, ..., N)$. However, the physical trajectories ($\vec{q}_i$ and $\vec{p}_i$) are $6N$ for an N-body system. $2N$ constraints are required to reduce the number of dimensions from $8N$ to physical trajectories $6N$ [42–46, 77–79].

$$\phi_i \approx 0 (i = 1, ..., 2N), \quad (11)$$

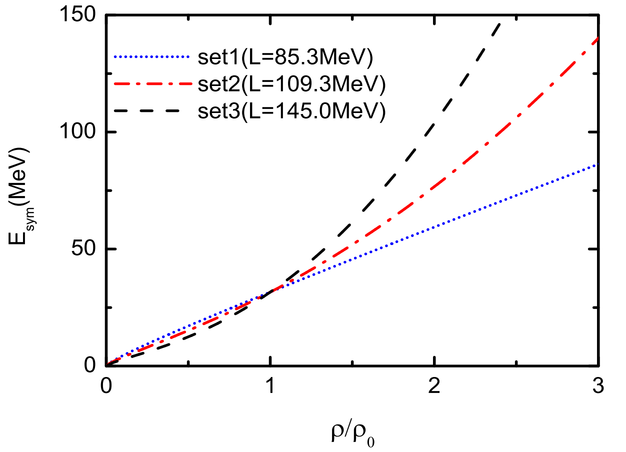

Fig. 2. (Color online) Symmetry energy as a function of the baryon density.

where $2N$ constraints satisfy the physical $6N$ phase space. The sign $\approx$ indicates Dirac's weak equality. The on-mass shell conditions can reduce the phase space from $8N$ to $7N$ dimensions.

$$\phi_i \equiv p_i^{*2} - M_i^{*2} = (p_i - V_i)^2 - (M_N - S_i)^2 = 0, \quad (12)$$

where $i = 1, ..., N$. The remaining N constraints are time fixation constraints. A simple choice of time fixation constraints that obey the worldline condition is written as [43, 77, 79, 80]

$$\phi_{i+N} \equiv \hat{a} \cdot (q_i - q_N) = 0, (i = 1, ..., N-1),$$
$$\phi_{2N} \equiv \hat{a} \cdot q_N - \tau = 0, \quad (13)$$

where $\hat{a} = (1, \vec{0})$ denotes a four-dimensional unit-vector [42–46, 77]. In a two-body center-of-mass system, $\hat{a}$ is defined as $p_{ij}^\mu / \sqrt{p_{ij}^2}$ with $p_{ij}^\mu = p_i^\mu + p_j^\mu$. We observe that only the constraint $i = 2N$ depends on τ. With the above $2N$ constraints, the number of dimensions from $8N$ will reduce to $6N$. These $2N$ constraints are conserved over time.

$$\frac{d\phi_i}{d\tau} = \frac{\partial \phi_i}{\partial \tau} + \sum_{k}^{2N} C_{ik}^{-1}\lambda_k = 0. \quad (14)$$

Because only the constraint $i = 2N$ depends on τ, λ is written as[77]

$$\lambda_i = -C_{2N,i}\frac{\partial \phi_{2N}}{\partial \tau}, (i = 1, ..., 2N-1), \quad (15)$$

with $C_{ij}^{-1} = [\phi_i, \phi_j]$. Following previous studies, the Hamiltonian of the N-body system was constructed as a linear combination of $2N-1$ constraints [77, 79, 80]:

$$H = \sum_{i=1}^{2N-1} \lambda_i(\tau)\phi_i, \quad (16)$$

Table 1. Parameter sets for RMF. The saturation density ρ_0 is set to 0.16 fm^{-3}. The binding energy of the saturation density is $E/A - M_N = -16$ MeV. The isoscalar-vector ω and isovector-vector ρ masses are fixed at their physical values, m_ω = 783 MeV and m_ρ = 763 MeV, respectively. The remaining meson masses, m_σ and m_δ, are set to be 550 and 500 MeV, respectively.

model	g_σ	g_ω	$g_2\,(fm^{-1})$	g_3	g_ρ	g_δ	K (MeV)	$E_{sym}(\rho_0)$ (MeV)	$L\,(\rho_0)$(MeV)	$M^*(\rho_0)/M_N$	K_{sym} (MeV)
set1	8.145	7.570	31.820	28.100	4.049	-	230	31.6	85.3	0.81	-15
set2	8.145	7.570	31.820	28.100	8.673	5.347	230	31.6	109.3	0.81	141
set3	8.145	7.570	31.820	28.100	11.768	7.752	230	31.6	145.0	0.81	391

Assuming $[\phi_i, \phi_j] = 0$, $\lambda_i = 0$ for $N+1 < i < 2N$ [77]. The equations of motion are then obtained as

$$\frac{dq_i}{d\tau} = [H, q_i] = \sum_j^N \lambda_j \frac{\partial \phi_j}{\partial p_i},$$

$$\frac{dp_i}{d\tau} = [H, p_i] = -\sum_j^N \lambda_j \frac{\partial \phi_j}{\partial q_i}, \qquad (17)$$

with the on-mass shell conditions (Eq. (12)) as inputs, the equations of motion can be obtained as

$$\dot{\vec{r}_i} = \frac{\vec{p}_i^*}{p_i^{*0}} + \sum_{j=1}^N \left(\frac{M_j^*}{p_j^{*0}} \frac{\partial M_j^*}{\partial \vec{p_i}} + z_j^{*\mu} \cdot \frac{\partial V_{j\mu}}{\partial \vec{p_i}}\right),$$

$$\dot{\vec{p}_i} = -\sum_{j=1}^N \left(\frac{M_j^*}{p_j^{*0}} \frac{\partial M_j^*}{\partial \vec{r_i}} + z_j^{*\mu} \cdot \frac{\partial V_{j\mu}}{\partial \vec{r_i}}\right), \qquad (18)$$

where $z_i^{*\mu} = p_i^{*\mu}/p_i^{*0}$ and $M_i^* = M_N - S_i$. The scalar potential S_i and vector potential $V_{i\mu}$ in RQMD are defined as

$$S_i = g_\sigma \sigma_i + g_\delta t_i \delta_i,$$

$$V_{i,\mu} = B_i g_\omega \omega_{i,\mu} + B_i t_i g_\rho b_{i,\mu}, \qquad (19)$$

where $t_i = 1$ for protons, and $t_i = -1$ for neutrons, and B_i is the baryon number of the ith particle. The meson field can be obtained from the RMF:

$$m_\sigma^2 \sigma_i + g_2 \sigma_i^2 + g_3 \sigma_i^3 = g_\sigma \rho_{S,i},$$

$$m_\omega^2 \omega_i^\mu = g_\omega J_i^\mu,$$

$$m_\delta^2 \delta_i = g_\delta(\rho_{Sp,i} - \rho_{Sn,i}) = g_\delta \rho_{S3,i},$$

$$m_\rho^2 b_i^\mu = g_\rho(\rho_{p,i} - \rho_{n,i}) = g_\rho R_i^\mu. \qquad (20)$$

Substituting Eq. (20) into Eq. (19), the scalar potential S_i and vector potential $V_{i\mu}$ of the nucleons can be obtained. For other hadrons, such as Δ resonances, similar to other transport models [34, 85], the Δ optical potential is estimated using the nucleon potentials and square of the Clebsch–Gordan coefficient. In the RQMD approach, the scalar density, isovector-scalar density, baryon current, and isovetor baryon current are expressed as

$$\rho_{S,i} = \sum_{j \neq i} \frac{M_j}{p_j^0} \rho_{ij}, \qquad \rho_{S3,i} = \sum_{j \neq i} t_j \frac{M_j}{p_j^0} \rho_{ij},$$

$$J_i^\mu = \sum_{j \neq i} B_j \frac{p_j^\mu}{p_j^0} \rho_{ij}, \quad R_i^\mu = \sum_{j \neq i} t_j B_j \frac{p_j^\mu}{p_j^0} \rho_{ij}. \qquad (21)$$

Because the difference between the numerical results obtained using the effective mass M_j^* and kinetic momentum $p_j^{\mu *}$ in the density and current and those obtained using a free mass $M_j = M_N = 938$ MeV and canonical momentum p_j^μ in the density and current is small, a free mass $M_j = M_N = 938$ MeV and canonical momentum p_j^μ have been used in the above density and current [45]. The interaction density ρ_{ij} is given by a Gaussian

$$\rho_{ij} = \frac{\gamma_{ij}}{(4\pi L_G)^{3/2}} \exp\left(\frac{q_{T,ij}^2}{4L_G}\right), \qquad (22)$$

where $q_{T,ij}^2 = q_{ij}^2 - [\frac{q_{ij,\sigma}(p_i^\sigma + p_j^\sigma)}{\sqrt{(p_i+p_j)^2}}]^2$ is a distance squared [77], and γ_{ij} is a Lorentz factor that ensures the correct normalization of the Gaussian [81] and is equal to $(p_i^0 + p_j^0)/(p_i + p_j)$ in the two-body center-of-mass frame. In this study, we set the Gaussian width to $L_G = 2.0 fm^2$.

C. Dispersion relation and production of pions

The Hamiltonian of the mesons is defined as [48, 82–84]

$$H_M = \sum_{i=1}^{N_M} [V_i^C + \omega(\vec{p_i}, \rho_i)], \qquad (23)$$

where V_i^C is the Coulomb potential, which is expressed as

$$V_i^C = \sum_{j=1}^{N_B} \frac{e_i e_j}{r_{ij}}, \qquad (24)$$

and N_M and N_B are the total numbers of mesons and baryons, including charged resonances, respectively. The pion potential in the medium, which contains isoscalar and isovector contributions, is defined as

$$\omega(\vec{p_i}, \rho_i) = \omega_{isoscalar}(\vec{p_i}, \rho_i) + C_\pi \tau_z \alpha (\rho/\rho_0)^{\gamma_\pi}, \quad (25)$$

where α denotes the isospin asymmetry parameter, the coefficient C_π is 36 MeV, the isospin quantity τ is 1, 0, and -1 for π^-, π^0, and π^+, respectively, and γ_π determines the isospin splitting of the pion potential and is set to two. In this study, the isoscalar part of pion potential $\omega_{isoscalar}$ was chosen as the Δ-hole model. The pion potential, which contains a pion branch (smaller value) and Δ-hole (larger value) branch, is defined as

$$\omega_{isoscalar}(\vec{p_i}, \rho_i) = S_\pi(\vec{p_i}, \rho_i)\omega_{\pi-like}(\vec{p_i}, \rho_i) + S_\Delta(\vec{p_i}, \rho_i)\omega_{\Delta-like}(\vec{p_i}, \rho_i). \qquad (26)$$

The probabilities of the pion and Δ-hole branches satisfy the following equation:

$$S_\pi(\vec{p_i},\rho_i) + S_\Delta(\vec{p_i},\rho_i) = 1. \tag{27}$$

The probability of both the pion and Δ-hole branches is defined as [84]

$$S(\vec{p_i},\rho_i) = \frac{1}{1-\partial\Pi(\omega)/\partial\omega^2}, \tag{28}$$

where ω denotes $\omega_{\pi-like}$ and $\omega_{\Delta-like}$. The eigenvalues of $\omega_{\pi-like}$ and $\omega_{\Delta-like}$ are generated from the pion dispersion relation

$$\omega^2 = \vec{p}_i^2 + m_\pi^2 + \Pi(\omega), \tag{29}$$

where Π denotes the pion self-energy. Including the short-range Δ-hole interaction, the pion self-energy is defined as

$$\Pi = \frac{\vec{p}_i^2\chi}{1-g'\chi}, \tag{30}$$

where m_π denotes pion mass. The Migdal parameter, g', was set to 0.6. χ is defined as

$$\chi = -\frac{8}{9}(\frac{f_\Delta}{m_\pi})^2\frac{\omega_\Delta\rho\hbar^3}{\omega_\Delta^2-\omega^2}\exp(-2\vec{p}_i^2/b^2), \tag{31}$$

where $\omega_\Delta = \sqrt{m_\Delta^2+\vec{p}_i^2} - M_N$ and m_Δ is the delta mass. In this study, the $\pi N\Delta$ coupling constant f_Δ was 2, and the cutoff factor b was $7m_\pi$.

Both the Coulomb and pion potentials contribute to the decay of resonances and the reabsorption of pions. For instance, the energy balance of Δ^0 in the decay of resonances and the reabsorption of pions can be written as

$$\sqrt{m_R^2+\vec{p}_R^2} = \sqrt{M_N^2+(\vec{p}_R-\vec{p}_\pi)^2} + \omega_\pi(\vec{p}_\pi,\rho) + V_\pi^C, \tag{32}$$

where $\vec{p}_R$ and $\vec{p}_\pi$ are the momenta of the resonances and pions, respectively, and m_R is the mass of the resonances. Because Δ^0 is uncharged, the Coulomb potential exists only for charged particles after the decay of Δ^0.

The pion is generated from direct nucleon-nucleon collision and decay of the resonances $\Delta(1232)$ and $N^*(1440)$. The reaction channels of the resonances and pions, which are the same as those in the LQMD model, are as follows [33, 48, 88, 89]:

$$NN \leftrightarrow N\Delta, \quad NN \leftrightarrow NN^*, \quad NN \leftrightarrow \Delta\Delta, \quad \Delta \leftrightarrow N\pi,$$
$$N^* \leftrightarrow N\pi, \quad NN \to NN\pi(s-state). \tag{33}$$

For the production of the $\Delta(1232)$ and $N^*(1440)$ resonances in nucleon-nucleon scattering, the parameterized cross-sections calculated using the one-boson exchange model were employed [91].

The decay width of $\Delta(1232)$ and $N^*(1440)$, which originates from the p-wave resonances, is momentum-dependent and is expressed as [91]

$$\Gamma(|\vec{p}|) = \frac{a_1|\vec{p}|^3}{(1+a_2|\vec{p}|^2)(a_3+|\vec{p}|^2)}\Gamma_0, \tag{34}$$

where $|\vec{p}|$ is the momentum of the created pion (in GeV/c). The parameters a_1, a_2, and a_3 were taken as 22.48 (17.22) c/GeV, 39.69 (39.69)c^2/GeV^2, and 0.04(0.09) GeV^2/c^2, respectively, for $\Delta(N^*)$. The bare decay width of $\Delta(N^*)$ was given by $\Gamma_0 = 0.12(0.2)$GeV. With the momentum-dependent decay width, the cross-section of pion-nucleon scattering has the Breit–Wigner form:

$$\sigma_{\pi N}(\sqrt{s}) = \sigma_{\max}(\frac{\vec{p}_m}{\vec{p}})^2\frac{0.25\Gamma^2(\vec{p})}{0.25\Gamma^2(\vec{p})+(\sqrt{s}-m_0)^2}, \tag{35}$$

where $\vec{p}$ and $\vec{p}_m$ are the three momenta of the pions at energies of $\sqrt{s}$ and m_0, respectively. The maximum cross-section $\sigma_{\max}$ of Δ and N^* resonances was obtained by fitting the total cross-sections of the experimental data in pion-nucleon scattering using the Breit–Wigner formula [92]. For instance, the maximum cross-section $\sigma_{\max}$ of Δ resonance was 200, 133.33, and 66.7 mb for $\pi^+p \to \Delta^{++}$ $(\pi^-n \to \Delta^-)$, $\pi^0p \to \Delta^+$ $(\pi^0n \to \Delta^0)$ and $\pi^-p \to \Delta^0$ $(\pi^+n \to \Delta^+)$, respectively [89].

Note that the threshold effect was neglected in this study. The threshold effect mainly refers to the Δ production threshold energy and incident energy of two colliding nucleons modified by the medium. The incident and threshold energies in the medium are defined as $\sqrt{s_{in}} = \sqrt{(E_1^*+\Sigma_1^0+E_2^*+\Sigma_2^0)^2-(\vec{\Sigma}_1+\vec{\Sigma}_2)^2}$ and $\sqrt{s_{th}} = \sqrt{(m_3^*+\Sigma_3^0+m_4^*+\Sigma_4^0)^2-(\vec{\Sigma}_3+\vec{\Sigma}_4)^2}$, respectively [85–87]. Because $\vec{\Sigma}_i = 0$ and $\vec{p}_i^* \simeq 0$ for static nuclear matter, the difference between the incident and threshold energies is $\sqrt{s_{in}} - \sqrt{s_{th}} = E_1^*+\Sigma_1^0+E_2^*+\Sigma_2^0-m_3^*-\Sigma_3^0-m_4^*-\Sigma_4^0$ [86]. The difference between the incident and threshold energies, which is isospin dependent, may result in an enhancement in the $nn \to p\Delta^-$ channel and suppression of the $pp \to n\Delta^{++}$ channel.

III. RESULTS AND DISCUSSIONS

The directed and elliptic flows were derived from the Fourier expansion of the azimuthal distribution:

$$\frac{dN}{d\phi}(y,p_T) = N_0[1+2V_1(y,p_T)\cos(\phi) +2V_2(y,p_T)\cos(2\phi)], \tag{36}$$

where the azimuthal angle of the emitted particle ϕ was measured from the reaction plane. $p_T = \sqrt{p_x^2+p_y^2}$ denotes the transverse momentum, and the directed flow V_1 and elliptic flow V_2 are expressed as follows:

$$V_1 \equiv <\cos(\phi)>=<\frac{p_x}{p_T}>,$$
$$V_2 \equiv <\cos(2\phi)>=<\frac{p_x^2-p_y^2}{p_T^2}>. \tag{37}$$

The directed flow provides information on the azimuthal anisotropy of the transverse emission. The elliptic flow describes the competition between the in-plane ($V_2 > 0$) and out-of-plane ($V_2 < 0$) emissions.

Firstly, the collective flows of LQMD.RMF in the ^{197}Au $+^{197}$ Au collisions have been investigated at an incident energy of 2.92 GeV/nucleon (the corresponding nucleon-nucleon center-of-mass energy is $\sqrt{S_{NN}} = 3$ GeV). The collective flows of LQMD with Skyrme interaction and without the momentum-dependent interaction have also been investigated. The Skyrme interaction of symmetric nuclear matter is taken to be the same as that of SLy6, with an incompressibility of $K = 230$ MeV at $\rho_0 = 0.16 fm^{-3}$[93, 94]. The symmetry energy of Skyrme interaction is defined as $E_{sym} = \frac{1}{3}\frac{\hbar^2}{2M_N}(\frac{3\pi^2\rho}{2})^{2/3} + \frac{C_{sym}}{3}(\rho/\rho_0)^{\gamma_s}$. When the C_{sym} and γ_s are set to be 38.6 MeV and 1.049, respectively, the symmetry energy and the slope parameter of symmetry energy are 31.6 MeV and 85.3 MeV, respectively. The symmetry energy and the slope parameter of symmetry energy of the Skyrme interaction are as same as those of set1. As shown in Fig.3, we have compared the collective flows of LQMD.RMF and LQMD with Skyrme interaction with recent experimental data from STAR Collaboration[90]. The collective follows of LQMD.RMF and LQMD with Skyrme interaction can describe the STAR data at an impact parameter b=4 fm and b=7 fm, respectively. The directed flows of LQMD.RMF are almost consistent with the STAR data across the entire rapidity. However, The directed flows of LQMD with Skyrme interaction are weaker than the STAR data across the entire rapidity. This result may be due to the fact that the value of directed transverse momentum with the Lorentz effect is larger than that without the Lorentz effect[43]. The elliptic flows of both LQMD.RMF and LQMD with Skyrme interaction are consistent with the STAR data at rapidities smaller than 0.5. However, the elliptic flows of both LQMD.RMF and LQMD with Skyrme interaction are weaker than the STAR data at large rapidity. At an incident energy of 2.92 GeV/nucleon, the LQMD.RMF can better describe the experimental data compared to the LQMD with Skyrme interaction and without the momentum-dependent interaction. Based on the above analyses, the RMF models have been implemented into the LQMD model successfully.

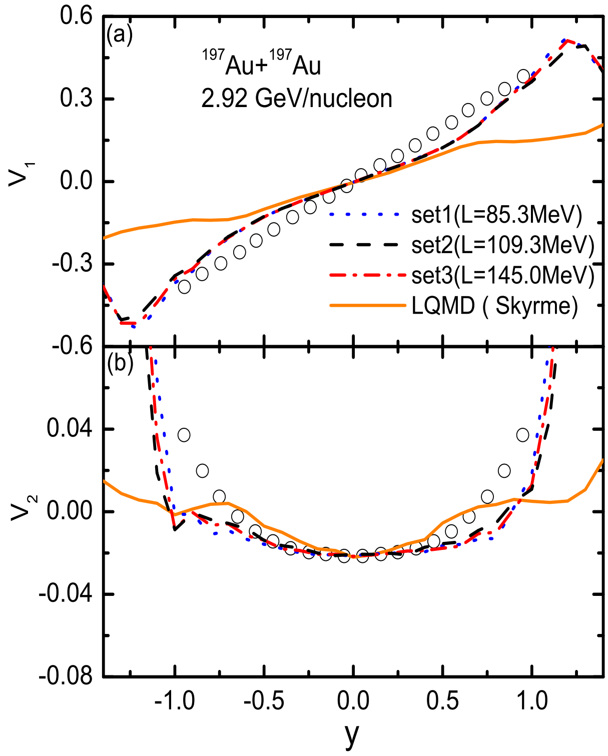

Fig. 3. (Color online) Rapidity distribution of the collective flows of free protons in the ^{197}Au $+^{197}$ Au reaction at an incident energy of 2.92 GeV/nucleon. The open circles correspond to the STAR data[90].

With this LQMD.RMF model, the ^{108}Sn $+^{112}$ Sn and ^{132}Sn $+^{124}$ Sn collisions in this study were investigated at an incident energy of 270 A MeV and an impact parameter b=3 fm. At an incident energy of 270 A MeV, the nuclear matter of the collision center can be compressed to densities approaching $2\rho_0$. In this dense region, collective flows, which reflect nucleon-nucleon interactions, can be used to extract the high-density behavior of the EOS [1, 44, 79, 88, 95]. The collective flows of free protons in the ^{108}Sn $+^{112}$ Sn and ^{132}Sn + ^{124}Sn collisions are shown in Fig. 4 and Fig. 5, respectively. It is reasonable that the maximum value of the directed flow V_1 was significantly larger than that of the elliptic flow V_2. In the same reaction system, the difference in the directed flows with various slopes of symmetry energy (set1, set2, and set3) was small. The difference in the elliptic flows with various slopes of symmetry energy was also small. To determine the relationship between the slope of the symmetry energy and the collective flow, we must process the collective flow data. The difference between neutron and proton directed flows emitted from heavy-ion collisions can be used to extract the density dependence of the symmetry energy [88, 95]. The difference between the neutron and proton directed flows is defined as $V_{1n} - V_{1p}$, as shown in Fig. 6. The trend and shape of the difference between the neutron and proton directed flows were similar to those of nonrelativistic LQMD [88]. For a given reaction system (nearly symmetric ^{108}Sn$+^{112}$Sn system or neutron-rich ^{132}Sn$+^{124}$Sn system), the absolute value of the difference between the neutron and proton directed flows with a soft symmetry energy was higher than that with a stiff symmetry energy. Interestingly, the stiffness of the symmetry energy can be reflected through the difference between the neutron and proton directed flows. The relationship between the curvature of the symmetry energy K_{sym} and the collective flows was then investigated. The difference between K_{sym} of set1 and K_{sym} of set3 was 406 MeV, which was significantly larger than the 59.7 MeV difference between L of set1 and L of set3. Although the curvature of the symmetry energy K_{sym} also affected the difference between the neutron and proton directed flows, because the nuclear matter of the collision center could only approach $2\rho_0$ at an incident energy of 270 A MeV, the effects of K_{sym} were not significant compared to the effects of the slope parameter L.

In addition to collective flows, the production of isospin exotic particles such as hyperons, kaons, and pions can also be used to extract the symmetry energy [22–28]. Because the thresholds of hyperons and kaons were not reached at the in-

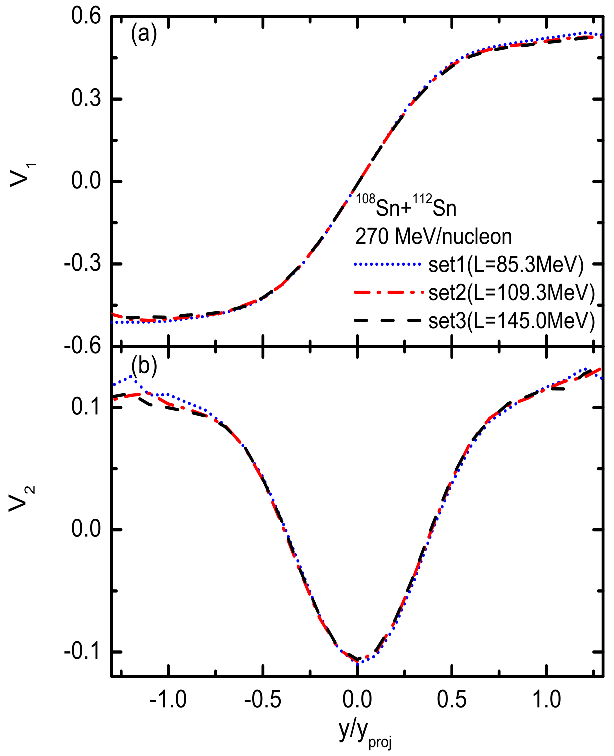

Fig. 4. (Color online) Rapidity distribution of the collective flows of free protons in the $^{108}\mathrm{Sn}+^{112}\mathrm{Sn}$ reaction at an incident energy of 270 MeV/nucleon.

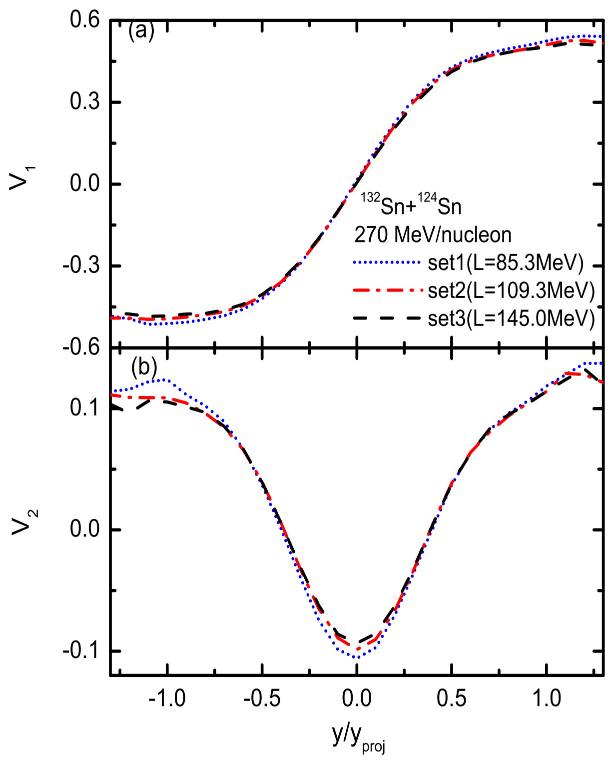

Fig. 5. (Color online) Rapidity distribution of the collective flows of free protons in the $^{132}\mathrm{Sn}+^{124}\mathrm{Sn}$ reaction at an incident energy of 270 MeV/nucleon.

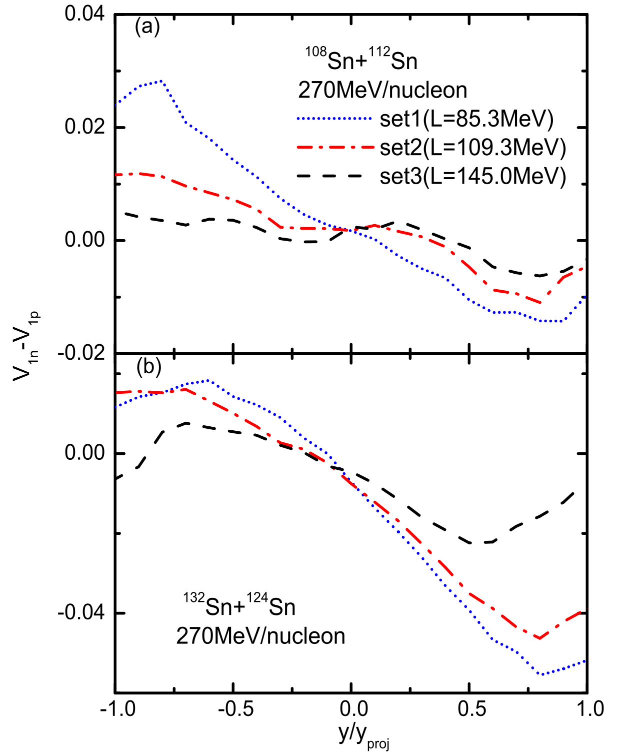

Fig. 6. (Color online) Difference between neutron and proton directed flows in the $^{108}\mathrm{Sn}+^{112}\mathrm{Sn}$ and $^{132}\mathrm{Sn}+^{124}\mathrm{Sn}$ reactions at an incident energy of 270 MeV/nucleon.

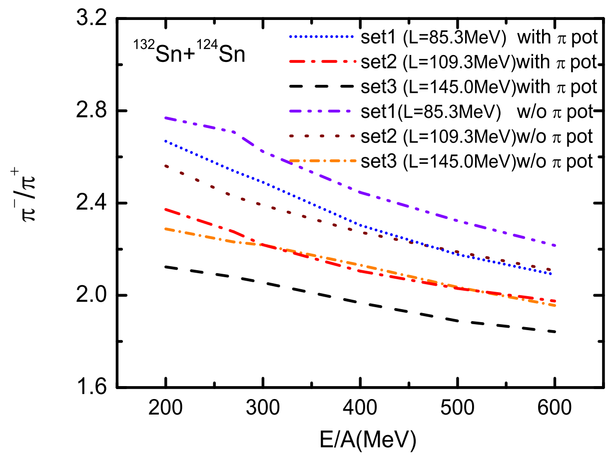

Fig. 7. (Color online) Ratio between the π^- and π^+ yields as a function of the incident energy in the $^{132}\mathrm{Sn}+^{124}\mathrm{Sn}$ reaction.

cident energies in this study, the isospin exotic particles were pions. First, we calculated the ratio between the π^- and π^+ yields of the neutron-rich $^{132}\mathrm{Sn}+^{124}\mathrm{Sn}$ system as a function of the collision energy at the impact parameter b=3 fm and $\theta_{cm} < 90°$. Because set1 had the softest symmetry energy, it had the highest neutron density. Consequently, there were more neutron-neutron scatterings in set1, resulting in the production of more Δ^- and π^-. As shown in Fig. 7, the π^-/π^+ ratio of set1 was the highest, and the π^-/π^+ ratio

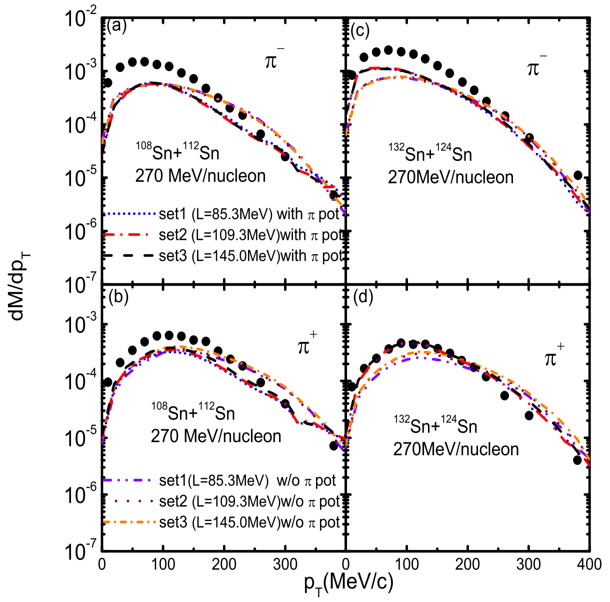

Fig. 8. (Color online) Transverse momentum spectra of pions as functions of transverse momentum at an incident energy of 270 MeV/nucleon. The left two panels [(a) and (b)] are the results of the $^{108}\text{Sn}+^{112}\text{Sn}$ reaction, and the right two panels [(c) and (d)] are the results of the $^{132}\text{Sn}+^{124}\text{Sn}$ reaction. The full circles correspond to the SπRIT data [40].

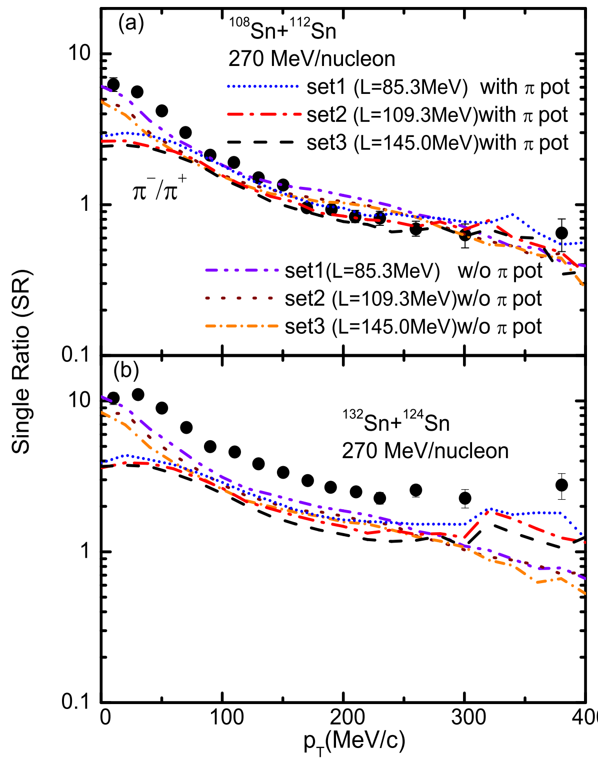

Fig. 9. (Color online) Single spectral ratios of pions as functions of transverse momentum for the $^{108}\text{Sn}+^{112}\text{Sn}$ and $^{132}\text{Sn}+^{124}\text{Sn}$ reactions at an incident energy of 270 MeV/nucleon. The full circles correspond to the SπRIT data [40].

of set2 was higher than that of set3. Specifically, at a collision energy of 270 MeV/nucleon, the π^-/π^+ ratio without (with) the π potential changed from 2.71 (2.54) to 2.23 (2.06) when the slope parameter was varied from $L=$85.3 to 145.0 MeV, that is, from set1 to set3. In other words, the π^-/π^+ ratio as a function of collision energy depends on the stiffness of the symmetry energy. This result was consistent with the predictions of most transport models [28, 39, 86]. When the π potential was considered, the interaction between π and the nucleon became attractive, resulting in a decrease in both π^- and π^+ via the absorbed channels $\pi N \to \Delta(1232)$ and $\Delta(1232)N \to NN$. However, with the π potential, because there were more neutrons in the neutron-rich $^{132}\text{Sn}+^{124}\text{Sn}$ system, π^- was more easily absorbed than π^+. Consequently, the π^-/π^+ ratio without the π potential was higher than that with the π potential in the same RMF model. Moreover, it is worth mentioning that the threshold effect neglected in this study may cause the π^-/π^+ ratio to be reversed. In other words, with the threshold effect [85–87], π^-/π^+ of set3 may be the highest, and the π^-/π^+ ratio of set2 may be higher than that of set1.

Next, the properties of π were predicted as a function of the transverse momentum. As shown in Fig. 8, the left and right panels are the transverse momentum spectra of pions for the nearly symmetric $^{108}\text{Sn}+^{112}\text{Sn}$ and neutron rich $^{132}\text{Sn}+^{124}\text{Sn}$ reactions at $\theta_{cm} < 90°$, respectively. In collisions between isotopes, π^+ is mainly generated from collisions between protons and π^- is mainly generated from collisions between neutrons. Theoretically, a stiffer symmetry energy would have a stronger repulsive force to push out neutrons and a stronger attractive force to squeeze protons, resulting in a decrease in the π^- yield and an increase in the π^+ yield, respectively. As shown in panels (b) and (d) of Fig. 8, the stiffer symmetry energy indeed led to larger transverse momentum spectra for π^+. However, the relationship between the transverse momentum spectra of π^- and the stiffness of the symmetry energy could not be directly explained. Compared with the stiffness of the symmetry energy, the π potential had a more significant impact on the transverse momentum spectrum of π. For the neutron-rich $^{132}\text{Sn}+^{124}\text{Sn}$ system, the transverse momentum spectra of both π^+ and π^- without the π potential were lower than those of the SπRIT data at $p_T \lesssim 200$ MeV. When the π potential was considered, the transverse momentum spectra of both π^+ and π^- increased at $p_T \lesssim 200$ MeV but decreased at $p_T \gtrsim 200$ MeV. Consequently, the transverse momentum spectra of π^+ with the π potential were almost consistent with the SπRIT data [40]; however, the transverse momentum spectra of π^- were still lower than the SπRIT data for the entire p_T domain. The lower transverse momentum spectra of π^- may be due to the absence of a threshold effect. The threshold effect, which was not considered in this study, may enhance the production of π^- [85–87].

Because a stiffer symmetry energy would have a stronger repulsive force to push out neutrons and a stronger attractive force to squeeze protons, resulting in a decrease in the π^- yield and an increase in the π^+ yield, respectively, the single

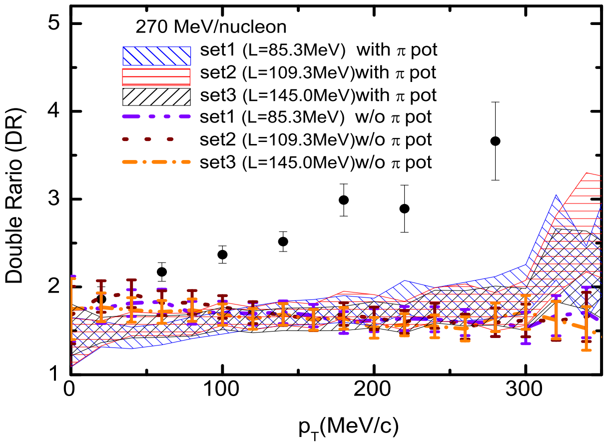

Fig. 10. (Color online) Double ratio of pions as a function of transverse momentum at an incident energy of 270 MeV/nucleon. The full circles correspond to the SπRIT data [40].

ratio SR$(\pi^-/\pi^+) = [dM(\pi^-)/dp_T]/[dM(\pi^+)/dp_T]$ may depend on the stiffness of the symmetry energy and the reaction system. As shown in Fig. 9, for both the nearly symmetric system and neutron-rich system, a softer symmetry energy led to a larger single ratio. In addition, for the same stiffness of the symmetry energy, because the neutron-neutron scattering of the neutron-rich $^{132}\mathrm{Sn}+^{124}\mathrm{Sn}$ system was greater than that of the nearly symmetric $^{108}\mathrm{Sn}+^{112}$ Sn system, the single ratio of the neutron-rich system was higher than that of the nearly symmetric system. It is worth mentioning that the single ratio of $^{108}\mathrm{Sn}+^{112}$ Sn was almost consistent with the experimental data. However, the single ratio of $^{132}\mathrm{Sn}+^{124}\mathrm{Sn}$ was lower than that of the experimental data for the entire p_T domain. This was because the transverse momentum spectra π^- of $^{132}\mathrm{Sn}+^{124}$ Sn were lower than the experimental data.

The double ratio between the neutron rich system and the nearly symmetric system DR(π^-/π^+) = SR$(\pi^-/\pi^+)_{132+124}$/SR$(\pi^-/\pi^+)_{108+112}$, which can cancel out most of the systematic errors caused by Coulomb and isoscalar interactions, was suggested to extract the properties of the symmetry energy [40]. However, because a lower symmetry energy will lead to a larger single ratio for both the nearly symmetric system and neutron-rich system, as shown in Fig. 10, it is still difficult to understand the dependence of the double ratio on the stiffness of symmetry energy. In addition, the double ratio without the π potential decreased slightly as a function of the transverse momentum, whereas it increased slightly as the transverse momentum increased when the π potential was considered. However, the increasing trend was still considerably weaker than that of the experimental results. The lower double ratio originated from the lower single ratio of the neutron-rich $^{132}\mathrm{Sn}+^{124}$ Sn system compared with the experimental data. The threshold effect may enhance the production of π^- [85–87] and the single ratio of the neutron-rich system, resulting in a larger double ratio.

IV. CONCLUSION

An RMF with various symmetry energies, namely set1, set2, and set3, was implemented in LQMD. The collective flows of the nearly symmetric $^{108}\mathrm{Sn}+^{112}$ Sn and neutron-rich $^{132}\mathrm{Sn}+^{124}$Sn systems were successfully generated from the LQMD.RMF. It has been observed that the directed flow V_1 was an order of magnitude larger than the elliptic flow V_2. However, the difference between the directed flows V_1 with various slopes of symmetry energy was small. The difference in the directed flows V_2 with various slopes of symmetry energy was also small. To explore the relationship between the collective flow and the stiffness of the symmetry energy, we defined the difference between neutron and proton directed flows $V_{1n} - V_{1p}$. For a given nearly symmetric system or neutron-rich system, the absolute value of $V_{1n} - V_{1p}$ increased with decreasing slope of the symmetry energy.

We also investigated the relationship between isospin exotic particles and the stiffness of the symmetry energy. The ratio between π^- yield and π^+ yield of the neutron-rich $^{132}\mathrm{Sn}+^{124}$ Sn system as a function of the collision energy increased with a decrease in the slope parameter of the symmetry energy. At an incident energy of 270 MeV/nucleon, the properties of π were predicted as a function of the transverse momentum. For the nearly symmetric $^{108}\mathrm{Sn}+^{112}$ Sn system, the single ratio of the nearly symmetric system was consistent with the experimental data. However, for the neutron-rich $^{132}\mathrm{Sn}+^{124}$ Sn system, the single ratio was lower than the experimental data, resulting in a double ratio lower than the experimental data. The π potential did not explain the lower transverse momentum spectra of π^- in the neutron-rich system. Considering the π potential, the double ratio increased slightly with increasing transverse momentum. However, the increasing trend was still considerably weaker than that observed in the experimental results. The single ratio of the neutron-rich system and the double ratio may be lower than the experimental data because of the lack of a threshold effect. The threshold effect, which can enhance the production of π^-, could be a candidate for enhancing the single ratio of the neutron-rich system to a double ratio. Moreover, because a softer symmetry energy led to a larger single ratio for both nearly symmetric and neutron-rich systems, the dependence of the double ratio on the stiffness of the symmetry energy was not significant. The sensitivity of the double ratio to the stiffness of the symmetry energy may also be due to the lack of a threshold effect. When the threshold effect is considered, the production of π^- in a neutron-rich system may be more significant than that in a nearly symmetric system. In the future, when collective flows occur, and the single ratio of the neutron-rich system and the double ratio of the LQMD.RMF are consistent with the experimental data, $V_{1n} - V_{1p}$ and the single ratio of the neutron-rich system π^-/π^+, which are sensitive to the stiffness of the symmetry energy, may be used to extract the slope parameter of the symmetry energy.

Appendix A: DETAILS of EQUATION OF MOTION

For numerical calculations, the equation of motion (Eq. (18)) must be written in the computed form. With Eq. (19) and $M_i^* = M_i - S_i = M_N - S_i$ as the inputs, the equation of motion (Eq. (18)) can be expanded as

$$\begin{aligned}\dot{\vec{r}}_i = \frac{\vec{p}_i^*}{p_i^{*0}} + \sum_{j\neq i}[D_{ij}\frac{\partial \rho_{ij}}{\partial \vec{p}_i} + D_{ji}\frac{\partial \rho_{ji}}{\partial \vec{p}_i} \\ +(D_j\frac{\partial f_i}{\partial \vec{p}_i} + A^{\mu}_{ji}\frac{\partial z_{i\mu}}{\partial \vec{p}_i})\rho_{ji} \\ +(D'_j t_i\frac{\partial f_i}{\partial \vec{p}_i} + A'^{\mu}_{ji}\frac{\partial z_{i\mu}}{\partial \vec{p}_i})\rho_{ji}],\end{aligned} \tag{A1}$$

$$\dot{\vec{p}} = -\sum_{j\neq i}[D_{ij}\frac{\partial \rho_{ij}}{\partial \vec{r}_i} + D_{ji}\frac{\partial \rho_{ji}}{\partial \vec{r}_i}], \tag{A2}$$

with

$$D_{ij} = D_i f_j + A^{\mu}_{ij} z_{j\mu} + D'_i t_j f_j + A'^{\mu}_{ij} z_{j\mu}, \tag{A3}$$

$$D_i = -g_\sigma \frac{M_i^*}{p_i^{*0}}\frac{\partial \sigma_i}{\partial \rho_{Si}}, \tag{A4}$$

$$A^{\mu}_{ij} = \frac{g_\omega^2}{m_\omega^2} B_i B_j z_i^{*\mu}, \tag{A5}$$

$$D'_i = -g_\delta t_i \frac{M_i^*}{p_i^{*0}}\frac{\partial \delta_i}{\partial \rho_{S3,i}}, \tag{A6}$$

$$A'^{\mu}_{ij} = \frac{g_\rho^2}{m_\rho^2} t_i t_j B_i B_j z_i^{*\mu}, \tag{A7}$$

where $z_i^\mu = p_i^\mu/p_i^0$ and $z_i^{*\mu} = p_i^{*\mu}/p_i^{*0}$. Based on Eq. (20), $\frac{\partial \sigma_i}{\partial \rho_{Si}}$ and $\frac{\partial \delta_i}{\partial \rho_{S3,i}}$ are obtained as follows:

$$\frac{\partial \sigma_i}{\partial \rho_{Si}} = \frac{g_\sigma}{m_\sigma^2 + 2g_2\sigma_i + 3g_3\sigma_i^2}, \quad \frac{\partial \delta_i}{\partial \rho_{S3,i}} = \frac{g_\delta}{m_\delta^2} \tag{A8}$$

In the two-body center-of-mass frame, ρ_{ij} equals ρ_{ji}. The squared distance $q^2_{T,ij}$ is reduced to $q^2_{T,ij} \equiv -\vec{r}^2_{ij} - \frac{(\vec{r}_{ij}\cdot\vec{p}_{ij})^2}{p^2_{ij}}$. In the actual calculation, we replace p_i^0 with $\sqrt{\vec{p}_i^2 + M_i^2}$ to save calculation time [79]. Thus, the partial derivative of density versus momentum and space can be written as

$$\begin{aligned}\frac{\partial \rho_{ij}}{\partial \vec{p}_i} = -\frac{\rho_{ij}}{2L}\frac{(\vec{r}_{ij}\cdot\vec{p}_{ij})\cdot\vec{r}_{ij}}{p^2_{ij}} + \rho_{ij}\frac{\gamma^2_{ij}\vec{\beta}_{ij}}{p_i^0 + p_j^0}[1 - \vec{\beta}_{ij}\frac{\vec{p}_i}{p_i^0}] \\ -\frac{\rho_{ij}}{2L}\frac{(\vec{r}_{ij}\cdot\vec{p}_{ij})^2}{p^4_{ij}}\{\vec{p}_{ij} - \frac{\vec{p}_i}{p_i^0}p^0_{ij}\},\end{aligned} \tag{A9}$$

$$\frac{\partial \rho_{ij}}{\partial \vec{r}_i} = -\frac{\rho_{ij}}{2L}[\vec{r}_{ij} + \frac{(\vec{r}_{ij}\cdot\vec{p}_{ij})\cdot\vec{p}_{ij}}{p^2_{ij}}], \tag{A10}$$

where $\vec{\beta}_{ij}$ is defined as $(\vec{p}_i + \vec{p}_j)/(p_i^0 + p_j^0)$. In addition, $\frac{\partial f_i}{\partial \vec{p}_i}$ can be written as $-\frac{M_i\vec{p}_i}{p_0^3}$. The partial derivative of the energy component of $\frac{\partial z_{i\mu}}{\partial \vec{p}_i}$ is zero and that of the momentum component of $\frac{\partial z_{i\mu}}{\partial \vec{p}_i}$ is written as $\frac{-p_i^2}{p_i^{03}}$. Using the above equations, the momentum and space of neutrons and protons are known.

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